

Dynamical Scenarios for $SU(2)_c$ Symmetric New Physics Interactions[†]

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Abstract

We study a picture of effective interactions among the W and Higgs bosons which is consistent with the precision tests at present energies, and at the same time allows for large observable New Physics (NP) effects in the bosonic sector. Toy dynamical models containing new heavy particles are used to indicate how such a picture could be created by integrating out the heavy particles. In these models custodial $SU(2)_c$ symmetry is realized at a certain large scale Λ_{NP} , either naturally or at a certain limit of a single coupling. In the derivation of the effective lagrangian we keep operators of dimension up to six. These operators involve only the known gauge boson and Higgs particles. They provide a valid description of NP up to an energy scale which is just below the mass of the new heavy degrees of freedom.

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Signals of New Physics (NP) beyond the Standard Model (SM) could appear in two different ways. The first possibility is that new particles associated with NP will be discovered at the future colliders. This spectacular case should appear at a future collider whose energy is larger than the typical mass of the new particles associated with NP. Standard arguments [1] suggest that this energy is in the TeV range, although it is not excluded that low lying states (like *e.g.* the lightest Higgs, the lightest supersymmetric particle, or new vector bosons) could be found earlier.

The second possibility is that all observable new particles are so heavy, that they cannot be directly produced by the future colliders. In such a case the only possibility left for an NP discovery is through precision tests probing departures from the SM predictions for processes involving usual particles. Such tests have already been performed for fermionic interactions at LEP1, SLC and in low energy experiments [2]. They indicate that the effective lagrangian describing NP should not include (at least at the few permille level) light fermion fields. The heavy quark sector has not yet been tested with such an accuracy though. For example, the $Z \rightarrow b\bar{b}$ width and the $b\bar{b}$ asymmetries just approach the percent level [3], while the consistency of top quark properties with the *SM* expectations are still to be confirmed at this level of accuracy [4]. On the other hand the bosonic sector (gauge boson self-interactions as well as the whole scalar sector) have not yet been directly tested at all.

It is generally expected (and hoped) that NP will partially reveal the mechanism responsible for generating the masses of the SM particles. Thus, it should not be surprising if NP signals first appear in sectors where mass generation plays an important role, like *e.g.* the heavy fermionic or the bosonic sector. From the agreement between the measured Z couplings to light fermions and the SM predictions, one concludes that if higher vector bosons exist and couple to light fermions with a standard electroweak strength, then their masses should lie at or above the 1 TeV range [5]. This limit lowers to a few hundred of GeV, if the new vector bosons couple very weakly to the light leptons and quarks [6]. Concerning the Higgs particle, no serious indication on its mass range (or to some extent even its existence) is yet provided, apart from a lower mass limit of about 60 GeV [7].

This leaves room for many possibilities for the appearance of NP, which include the cases demanding that the Higgs boson is replaced by a strongly interacting sector at the TeV scale [1, 8]. Thus in the present paper, we concentrate in a framework where all new degrees of freedom are very heavy, and the interactions among the known particles up to a scale $\Lambda_{NP} \gtrsim 1$ TeV, are described by an effective lagrangian satisfying $SU(2) \times U(1)$ gauge symmetry and consisting of the SM lagrangian, plus additional NP operators involving only the known gauge boson and Higgs fields. The order of magnitude of the contributions of these NP operators is determined by the scale Λ_{NP} . To lowest nontrivial order we restrict to NP operators of dimension up to six [9].

These NP operators induce both direct and indirect contributions to the 2-point gauge boson functions (γ , Z , W self-energies). The direct ones appear at the tree level and they are strongly suppressed due to the present precision measurements. The indirect contributions though, induced by the use of the effective interactions inside loops, must be handled with care in order to be meaningful. Reliable constraints can then only be

obtained from effective lagrangians which lead to a decent high energy behaviour and do not produce results strongly depending on the scale Λ_{NP} ; *i.e.* on the details of the regime above or close to the NP threshold. Such a behaviour can best be achieved if the NP operators satisfy $SU(2) \times U(1)$ gauge invariance [10], which is of course mandatory for any decent theory. In this respect it is interesting to mention that it is always possible to write any interaction respecting electromagnetic gauge invariance, in a $SU(2) \times U(1)$ gauge invariant form, provided no constraint on the dimensionality of the allowed operators is imposed [11, 12]. So $SU(2) \times U(1)$ gauge invariance alone does not restrict the kind of possible interactions, since the implied improvement of the high energy behaviour can always be achieved either by the additional multi-boson vertices or by vertices involving Higgses. However, if $\Lambda_{NP} \gg v$ (v is the electroweak scale), restrictions will be provided by the dimension of the effective operators.

In addition to $SU(2) \times U(1)$ gauge invariance, in the present paper we assume that NP satisfies CP conservation and restrict to purely bosonic operators of dimension up to six. This assumption is motivated by remarking that the fermionic contribution to the $SU(2) \times U(1)$ currents may always be eliminated by using the gauge boson equations of motion [9], while the remaining fermionic terms should somehow be related to the mechanism responsible for the spontaneous symmetry breaking. Because of this, their strength may be similar to the Yukawa type fermionic couplings which are known to be small, except for the top quark case. Thus, following [13] we disregard all fermionic operators and describe NP in terms of boson operators only. The equations of motion are then used only so far they relate bosonic operators among themselves. The a priori very long list of these operators reduces to only 11 independent ones.

Four of these operators may be written as,

$$\overline{\mathcal{O}}_{DW} = 4 \langle ([D_\mu, W^{\mu\rho}])([D^\nu, W_{\nu\rho}]) \rangle, \quad (1)$$

$$\mathcal{O}_{DB} = (\partial_\mu B_{\nu\rho})(\partial^\mu B^{\nu\rho}), \quad (2)$$

$$\mathcal{O}_{BW} = \Phi^\dagger B_{\mu\nu} W^{\mu\nu} \Phi, \quad (3)$$

$$\mathcal{O}_{\Phi 1} = (D_\mu \Phi)^\dagger \Phi \Phi^\dagger (D^\mu \Phi). \quad (4)$$

In addition to these we also give for later convenience the redundant operator¹

$$\mathcal{O}_{\Phi 4} = (\Phi^\dagger \Phi)(D_\mu \Phi)^\dagger (D^\mu \Phi). \quad (5)$$

All these operators contribute directly to the gauge boson 2-point functions, and according to the precision tests, they should have reduced couplings. In the preceding equations,

¹The operator $\mathcal{O}_{\Phi 4}$ of (5) was first introduced in [9]. It was then omitted in [13] because partial integration of $\mathcal{O}_{\Phi 2}$ (see (9)), leads to $\mathcal{O}_{\Phi 2} + 8\mathcal{O}_{\Phi 4} + 4(\Phi^\dagger \Phi) \{ (D_\mu D^\mu \Phi)^\dagger \cdot \Phi + \Phi^\dagger \cdot (D_\mu D^\mu \Phi) \} = 0$, which combined with the Higgs equations of motion (where the Yukawa contributions are neglected) implies that $\mathcal{O}_{\Phi 4}$ is equivalent to a combination of purely Higgs operators without any gauge boson couplings. Note that the substitution of $\mathcal{O}_{\Phi 4}$ by these Higgs fields implies a renormalization of the parameters already existing in SM, which induces the necessary changes to the gauge couplings explicitly modified by the presence of $\mathcal{O}_{\Phi 4}$.

the definitions

$$\Phi = \left(\begin{array}{c} \phi^+ \\ \frac{1}{\sqrt{2}}(v + H + i\phi^0) \end{array} \right) \quad , \quad (6)$$

$$\begin{aligned} D_\mu &= (\partial_\mu + i g_1 Y B_\mu + i g_2 W_\mu) \quad , \\ W_\mu &= \vec{W}_\mu \cdot \frac{\vec{\tau}}{2} \quad , \quad W_{\mu\nu} = \vec{W}_{\mu\nu} \cdot \frac{\vec{\tau}}{2} \quad , \end{aligned} \quad (7)$$

are used, where Y gives the hypercharge of the field to which the covariant derivative acts, v is vacuum expectation value of the Higgs field and $\langle A \rangle \equiv Tr A$. For studying the $SU(2)_c$ transformations of the NP operators, we quote the expression for the Higgs field

$$\hat{U} = (\tilde{\Phi} \quad , \quad \Phi) \quad , \quad (8)$$

where $\tilde{\Phi} = i\tau_2 \Phi^*$.

Another two NP operators, constructed with scalar fields only, are

$$\mathcal{O}_{\Phi 2} = (\partial_\mu \langle \hat{U} \hat{U}^\dagger \rangle) (\partial^\mu \langle \hat{U} \hat{U}^\dagger \rangle) \quad , \quad (9)$$

$$\mathcal{O}_{\Phi 3} = \langle \hat{U} \hat{U}^\dagger \rangle^3 \quad . \quad (10)$$

We call them "superblind", since they do not give observable contributions to the present precision measurements even at the one-loop level.

Finally the last five operators

$$\mathcal{O}_W = \frac{1}{3!} \left(\vec{W}_\mu{}^\nu \times \vec{W}_\nu{}^\lambda \right) \cdot \vec{W}_\lambda{}^\mu = -\frac{2i}{3} \langle W^{\nu\lambda} W_{\lambda\mu} W^\mu{}_\nu \rangle \quad , \quad (11)$$

$$\hat{\mathcal{O}}_{UW} = \frac{1}{2} \langle \hat{U} \hat{U}^\dagger \rangle \langle W^{\mu\nu} W_{\mu\nu} \rangle \quad , \quad (12)$$

$$\hat{\mathcal{O}}_{UB} = \langle \hat{U} \hat{U}^\dagger \rangle B^{\mu\nu} B_{\mu\nu} \quad , \quad (13)$$

$$\mathcal{O}_{W\Phi} = 2 (D_\mu \Phi)^\dagger W^{\mu\nu} (D_\nu \Phi) \quad , \quad (14)$$

$$\mathcal{O}_{B\Phi} = (D_\mu \Phi)^\dagger B^{\mu\nu} (D_\nu \Phi) \quad , \quad (15)$$

constitute the most interesting set since the present precision experiments still allow them to be appreciable.

Note that in this list of independent operators, we have used $\overline{\mathcal{O}}_{DW}$ given in (1), instead of the one usually called \mathcal{O}_{DW} defined as

$$\mathcal{O}_{DW} = 2 \langle ([D_\mu, W_{\nu\rho}]) ([D^\mu, W^{\nu\rho}]) \rangle \quad , \quad (16)$$

and related to \mathcal{O}_W and $\overline{\mathcal{O}}_{DW}$ by the identity

$$\mathcal{O}_{DW} = 12 g_2 \mathcal{O}_W + \overline{\mathcal{O}}_{DW} \quad . \quad (17)$$

The reason is the following. From the expressions of these operators given in (11, 1, 16) we observe that $\overline{\mathcal{O}}_{DW}$ and \mathcal{O}_{DW} generate the same contribution to the 2-point gauge function at the tree level. The difference between these operators first arises in the triple gauge vertices. A study reveals that the triple gauge vertex in $\overline{\mathcal{O}}_{DW}$ has the usual Yang-Mills form (apart from additional d'Alembertian derivatives), whereas \mathcal{O}_{DW} includes in addition the genuinely anomalous 3-gauge quadruple coupling described by \mathcal{O}_W [14]. This later coupling does not appear in $\overline{\mathcal{O}}_{DW}$.

The five operators appearing in (11-16) are called "blind" [10], since they contribute to the quantities measured in the precision experiments at LEP1, only at the one-loop level². As one could imagine from a rough estimate based on the $(\alpha/4\pi)$ loop factors, these measurements only produce mild indirect constraints on the couplings of these operators. This is what comes out when effective operators are treated one by one. If several operators are simultaneously considered though, cancellation effects appear and even these mild constraints essentially vanish [13]. Because of this, present measurements allow the couplings of the blind operators to be considerable, which in turn makes the operators themselves interesting.

The energy domain where such operators can appear as local, has also been discussed on the basis of unitarity considerations [15, 16]. Locality, together with high dimensionality, will inevitably produce amplitudes which approach the unitarity limit at a sufficiently high scale. This phenomenon is well known from the old Fermi current-current interaction, which is also a $dim = 6$ local operator. As in the Fermi theory, these unitarity constraints give indications not only on the domain of validity of the effective theory, but also on the threshold of the related NP. Thus, depending on the type of operators, at sufficiently high energies either strong interactions will appear involving gauge bosons of various helicities, or new degrees of freedom will be excited³. The scale where this happens can be identified as Λ_{NP} . Relations between the NP coupling constants and the saturation scales Λ_{NP} have been established for each operator [15, 16]. Assuming that Λ_{NP} is of the order of 1 TeV, one then gets "unitarity" upper bounds for these "anomalous" couplings, which are as stringent as those obtained indirectly from the LEP1 measurements. Alternatively, if an upper bound on any of these couplings is experimentally established, then the unitarity relations may be used to give a lower bound on the threshold of the related NP.

To summarize the present situation concerning the five "blind" operators, we state that LEP1 and unitarity do not strongly constrain them. In principle much stronger constraints will be obtained from direct measurements on these operators. One has of course to check if this is achievable in the contemplated future experiments. In this respect we note that at LEP2 [17], couplings of the above "blind" operators at the level of at least $O(0.1)$ are expected to be visible. Subsequently, LHC [18, 19], and NLC [20, 21] should allow to reach the $O(0.01)$ level, or even the permille level, for such NP couplings. In addition, laser back scattering experiments may turn out to be especially stringent for studying *e.g.* the $H\gamma\gamma$ anomalous (and normal) coupling, in case a light Higgs boson

²For the peculiarities of $\widehat{\mathcal{O}}_{UW}$ and $\widehat{\mathcal{O}}_{UB}$ see [15] and below.

³In fact this later possibility would be analogous to the Fermi case, where the W excitation destroyed the locality of the current interaction.

could be produced through the $\gamma\gamma$ collisions [22, 23]. So finally the same level of accuracy should be reached for the bosonic sectors at NLC, as for the fermionic sector at LEP1. It makes therefore sense to discuss the NP effects in these various sectors and from their comparison, to try to infer the NP properties.

The fact that no NP effect has been seen in the fermionic sector at LEP1, and the possibility that such an effect could be seen in the bosonic sectors, should not be considered as an "unnatural" situation, but rather as a remarkable signal of the NP properties. Like *e.g.* the presence of a certain symmetry. An example of such symmetry is custodial $SU(2)_c$. Cumulating the information that lepton and quark couplings seem standard, and that neither W-B mixing nor $\Delta\rho$ effects in the self-energy ϵ parameters appear in nature, provides ample motivation for the possible importance of this global symmetry for NP [24, 18, 15]. This motivation is further supported by the expectation that NP may be related to the scalar sector of the electroweak interactions, which is known to preserve $SU(2)_c$ in SM, whereas the B_μ and fermionic couplings violate it.

An $SU(2)_c$ preserving NP effective lagrangian should include neither ordinary fermions nor the B_μ fields. It can only involve the W_μ and Higgs fields⁴. If we assume that NP somehow respects global $SU(2)_c$ symmetry, then the above set of 11 operators is restricted to only 5 ones namely, the two superblinds \mathcal{O}_{Φ_2} , \mathcal{O}_{Φ_3} , the two operators \mathcal{O}_W and $\overline{\mathcal{O}}_{DW}$ involving W_μ only, and the operator $\widehat{\mathcal{O}}_{UW}$ involving both W_μ and Higgs fields.

At this point it is worthwhile to offer examples of dynamical models for NP which at the effective lagrangian level have to some extent the $SU(2)_c$ symmetry mentioned above. Thus in model A below, the effective NP interaction at the large scale Λ_{NP} obeys global $SU(2)_c$ invariance in the limit that a certain coupling called $g_{\psi 2}$ vanishes, while in models B and B' this symmetry is realized independently of the actual values of the couplings. At a lower scale the couplings of these operators will of course run according to the rules implied by the SM lagrangian. Consequently smaller contributions from $SU(2)_c$ violating operators will also be generated at such lower scales.

Model A:

In this model, we assume that NP is determined by a complex scalar field Ψ which under the $SU(2) \times U(1)$ gauge group has isospin I and vanishing hypercharge. The associated Ψ particles are assumed to have a large mass $M = \Lambda_{NP}$, and possibly also some hyper-colour \tilde{N}_c . The basic renormalizable lagrangian will then be the sum

$$\mathcal{L} = \mathcal{L}_{SM} + \mathcal{L}_\psi \quad , \quad (18)$$

of the SM lagrangian

$$\begin{aligned} \mathcal{L}_{SM} = & -\frac{1}{2}\langle W_{\mu\nu}W^{\mu\nu} \rangle - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} + \frac{1}{2}\langle D_\mu \widehat{U} D^\mu \widehat{U}^\dagger \rangle \\ & - \frac{M_H^2}{8v^2} \left(\langle \widehat{U} \widehat{U}^\dagger \rangle - v^2 \right)^2 + \text{fermionic terms} \quad , \end{aligned} \quad (19)$$

⁴ This use of $SU(2)_c$ symmetry is somewhat stronger than the one where it is applied only after the B_μ field has been removed ($g_1 \rightarrow 0$) [24].

and the lagrangian describing the new degrees of freedom

$$\mathcal{L}_\psi = (D_\mu \Psi)^\dagger (D^\mu \Psi) - \Psi^\dagger \left(\Lambda_{NP}^2 - g_{\psi 1} \langle \hat{U} \hat{U}^\dagger \rangle - g_{\psi 2} \hat{U} \tau_3 \hat{U}^\dagger \right) \Psi \quad , \quad (20)$$

where the definition (8) has been used. The tensorial representation of the isospin= I field Ψ is implicitly used, so that *e.g.* in the last term of (20) one of the isospin indices of Ψ and Ψ^\dagger are dotted to $\hat{U} \tau_3 \hat{U}^\dagger$, while the remaining indices are dotted among themselves. This term violates $SU(2)_c$. We note that (20) gives, apart from an irrelevant $(\Psi^\dagger \Psi)^2$ term, the most general renormalizable interaction in consistence with $SU(2) \times U(1)$ gauge invariance. The vanishing hypercharge of Ψ implies that no B_μ interactions appear in (20).

The standard techniques [25] may now be used to calculate the effective lagrangian describing the electroweak interactions at a scale μ just below Λ_{NP} . This is obtained by integrating out at the one loop order the heavy degrees of freedom described by the field Ψ . Thus, at the scale just below Λ_{NP} , the effective electroweak interaction among the usual SM particles, is obtained by employing the Seeley-DeWitt expansion of the relevant determinant⁵ [26]. It is given by

$$\mathcal{L}_{eff} = \mathcal{L}_{SM} + \mathcal{L}_{NP} \quad , \quad (21)$$

$$\begin{aligned} \mathcal{L}_{NP} = & \frac{(2I+1)\widetilde{N}_c}{(4\pi)^2} \left\{ - g_{\psi 1} \Lambda_{NP}^2 \left(\frac{1}{\epsilon} + 1 \right) \langle \hat{U} \hat{U}^\dagger \rangle + \frac{1}{2\epsilon} \left(g_{\psi 1}^2 + g_{\psi 2}^2 \frac{I(I+1)}{3} \right) \langle \hat{U} \hat{U}^\dagger \rangle^2 \right. \\ & - \frac{g_2^2 I(I+1)}{18\epsilon} \langle W_{\mu\nu} W^{\mu\nu} \rangle + \frac{1}{6\Lambda_{NP}^2} (g_{\psi 1}^3 + g_{\psi 1} g_{\psi 2}^2 I(I+1)) \langle \hat{U} \hat{U}^\dagger \rangle^3 \\ & + \frac{1}{12\Lambda_{NP}^2} \left[\left(g_{\psi 1}^2 + g_{\psi 2}^2 \frac{I(I+1)}{3} \right) \mathcal{O}_{\Phi 2} + g_{\psi 2}^2 \frac{16I(I+1)}{3} (\mathcal{O}_{\Phi 4} - \mathcal{O}_{\Phi 1}) \right] \\ & \left. - \frac{g_2^2 I(I+1)}{90\Lambda_{NP}^2} \left[10g_{\psi 1} \hat{\mathcal{O}}_{UW} + g_2 \mathcal{O}_W + \frac{1}{4} \overline{\mathcal{O}}_{DW} \right] \right\} \quad , \quad (22) \end{aligned}$$

where the definitions (1,4,5,8,9,11,12) have been used, and $\epsilon = (2 - n/2)$ is the usual dimensional regularization parameter for the ultraviolet divergences.

The first three terms in (22) renormalize $SU(2)_c$ invariant quantities already existing in the SM lagrangian (19), while the remaining ones create contributions from the five possible $dim = 6$, $SU(2)_c$ conserving operators $\mathcal{O}_{\Phi 2}$, $\mathcal{O}_{\Phi 3}$, \mathcal{O}_W , $\overline{\mathcal{O}}_{DW}$ and $\hat{\mathcal{O}}_{UW}$. Finally $\mathcal{O}_{\Phi 4} - \mathcal{O}_{\Phi 1}$ gives the only $SU(2)_c$ violating but B_μ independent contribution, which is generated by the $g_{\psi 2}$ coupling. The fact that no B_μ couplings appear in (22) is a direct consequence of the vanishing hypercharge of the Ψ field.

The term $W_{\mu\nu} W^{\mu\nu}$ in (22) induces a non observable renormalization of W_μ , while the purely Higgs operators renormalize the vacuum expectation value and the mass of the

⁵ This result has been also checked by an explicit calculation of the self energy and triangle loops.

Higgs field, producing in addition anomalous self-interactions. After this renormalization is done we are lead to

$$\mathcal{L}_{NP} = \frac{2\tilde{d}}{v^2} \hat{\mathcal{O}}_{UW} + \frac{4\tilde{\lambda}_W}{g_2 v^2} \mathcal{O}_W + \frac{4\tilde{x}_{\overline{DW}}}{g_2^2 v^2} \overline{\mathcal{O}}_{DW} + \frac{g_\phi}{\Lambda_{NP}^2} (\mathcal{O}_{\Phi 4} - \mathcal{O}_{\Phi 1}) + (\text{Higgs self-interactions}) . \quad (23)$$

Using (21,22) and $v = 2M_W/g_2$, we get

$$\tilde{\lambda}_W = - \frac{(2I+1)I(I+1)\tilde{N}_c}{90 (4\pi)^2} \frac{g_2^2 M_W^2}{\Lambda_{NP}^2} , \quad (24)$$

$$\frac{\tilde{x}_{\overline{DW}}}{\tilde{\lambda}_W} = \frac{1}{4} , \quad (25)$$

$$\frac{\tilde{d}}{\tilde{\lambda}_W} = \frac{20 g_{\psi 1}}{g_2^2} . \quad (26)$$

Eliminating now the trivial contribution to the W kinetic energy from $\hat{\mathcal{O}}_{UW}$ by renormalizing the W_μ field and g_2 as in [15] we get

$$\mathcal{L}_{NP} = d\mathcal{O}_{UW} + \frac{\lambda_W g_2}{M_W^2} \mathcal{O}_W + \frac{x_{\overline{DW}}}{M_W^2} \overline{\mathcal{O}}_{DW} + \frac{g_\phi}{\Lambda_{NP}^2} (\mathcal{O}_{\Phi 4} - \mathcal{O}_{\Phi 1}) + (\text{Higgs self-interactions}), \quad (27)$$

where [15]

$$\mathcal{O}_{UW} = \frac{1}{v^2} \langle \hat{U} \hat{U}^\dagger - \frac{v^2}{2} \rangle \langle W_{\mu\nu} W^{\mu\nu} \rangle , \quad (28)$$

and

$$d = \frac{\tilde{d}}{1 - 2\tilde{d}} , \quad (29)$$

$$\lambda_W = \frac{\tilde{\lambda}_W}{(1 - 2\tilde{d})^2} , \quad x_{\overline{DW}} = \frac{\tilde{x}_{\overline{DW}}}{(1 - 2\tilde{d})^2} , \quad (30)$$

satisfying again

$$\frac{x_{\overline{DW}}}{\lambda_W} = \frac{1}{4} . \quad (31)$$

Note that the difference between the tilded couplings in (24-26) and those without it in (28-30), arises from the respective use of $\hat{\mathcal{O}}_{UW}$ and \mathcal{O}_{UW} in \mathcal{L}_{NP} ; compare (23, 27). To lowest order in the electroweak factor $g_2^2/(4\pi)^2$ this difference vanishes. The magnitude of these couplings is controlled by $g_2^2/(4\pi)^2$, the $dim = 6$ scale ratio M_W^2/Λ_{NP}^2 and the isospin and hyper-colour coefficients. Note the ratio of 4/1 in favour of \mathcal{O}_W over $\overline{\mathcal{O}}_{DW}$, and the negative signs of both couplings. Since $\overline{\mathcal{O}}_{DW}$ is already strongly constrained by LEP1 [13], this result provides only little chance for the observability of λ_W .

One also sees that the three Higgs involving $SU(2)_c$ invariant operators $\mathcal{O}_{\Phi 2}$, $\mathcal{O}_{\Phi 3}$ and \mathcal{O}_{UW} owe their presence to the basic coupling $g_{\psi 1}$ in (20). Most interesting is the

operator \mathcal{O}_{UW} whose physical consequences have been studied in [19, 21, 23]. Contrary to the case for the operators \mathcal{O}_W and $\overline{\mathcal{O}}_{DW}$, the coupling of \mathcal{O}_{UW} could not be predicted even if Λ_{NP} were known, because $g_{\psi 1}$ is unknown. Nevertheless it is interesting to remark that according to (26,29,30), d may be considerably larger than λ_W if $g_{\psi 1}$ is comparable to the electroweak coupling g_2^2 . Finally the only $SU(2)_c$ violating contribution in this model arises from $\mathcal{O}_{\Phi 4} - \mathcal{O}_{\Phi 1}$ and is determined by the other unknown coupling $g_{\psi 2}$. The single direct contribution of this operator is to Π_{WW} , which in turn implies that the only precision measurement parameter which is sensitive to it is $\epsilon_1 = \Delta\rho$. Present precision measurements constrain $g_{\psi 2}$ to be very small.

Model B:

We now try a second toy model which naturally generates only $SU(2)_c$ invariant effective interactions. In this model NP is given by a complex fermion field F whose left and right components have both vanishing hypercharge, isospin I and hyper-colour \widetilde{N}_c . The associated particles again have a large gauge invariant mass $M_f = \Lambda_{NP}$. The basic lagrangian is now

$$\mathcal{L} = \mathcal{L}_{SM} + \mathcal{L}_F \quad , \quad (32)$$

where (19) is used and the most general $SU(2) \times U(1)$ gauge invariant and renormalizable lagrangian for NP is

$$\mathcal{L}_F = i\overline{F}(\not{\partial} + ig_2\overline{W} \cdot \overrightarrow{t})F - \Lambda_{NP}\overline{F}F \quad , \quad (33)$$

with \overrightarrow{t} denoting the isospin I matrices.

Integrating the fermion loop as before [26] and renormalizing appropriately W_μ , we get at the scale Λ_{NP}

$$\mathcal{L}_{NP} = \frac{(2I+1)I(I+1)\widetilde{N}_c}{(4\pi)^2} \frac{g_2^2}{45\Lambda_{NP}^2} [g_2\mathcal{O}_W - \overline{\mathcal{O}}_{DW}] \quad . \quad (34)$$

Comparing with (23, 27-30), we conclude that there is no \mathcal{O}_{UW} contribution in this model, and

$$\lambda_W = \frac{(2I+1)I(I+1)\widetilde{N}_c}{45(4\pi)^2} \frac{g_2^2 M_W^2}{\Lambda_{NP}^2} \quad , \quad (35)$$

$$x_{\overline{DW}} = -\lambda_W \quad . \quad (36)$$

Several other nontrivial features can also be noticed. Thus, the sign of λ_W is now positive, opposite to that of $x_{\overline{DW}}$ and opposite to λ_W in model A. Moreover \mathcal{O}_W and $\overline{\mathcal{O}}_{DW}$ are now generated with equal strength. As before the actual magnitude of the couplings is controlled by the isospin and hyper-colour factors.

At this stage no NP couplings to the Higgs field are generated in model B. In order to get them the model is extended to

Model B' :

To the preceding NP spectrum we just add a heavy scalar isoscalar field S^0 . The most general gauge invariant and renormalizable lagrangian is now given by (32), with (33) with \mathcal{L}_F given by

$$\mathcal{L}_F = i\bar{F}(\not{\partial} + ig_2 \not{W})F - M_f \bar{F}F + f_f S^0 \bar{F}F + \frac{f_\phi}{2} M_s S^0 \langle \hat{U} \hat{U}^\dagger \rangle + \frac{1}{2} (\partial_\mu S^0) (\partial^\mu S^0) - \frac{M_s^2}{2} (S^0)^2, \quad (37)$$

where⁶ the S^0 and F masses are assumed to be similar; *i.e.* $M_s \sim M_f \equiv \Lambda_{NP}$. So essentially only two additional parameters have been introduced in this extension, namely the dimensionless couplings f_f and f_ϕ . A priori, they are also of order 1.

By integrating the heavy fermion loop as before [26], we obtain the effective lagrangian, which as far as the \mathcal{O}_W and \mathcal{O}_{DW} contributions are concerned, is the same as the one given in (34). In addition to it though, the heavy fermion loop also creates an effective $WW S$ coupling of the form

$$\mathcal{L}_{SWW} = \frac{x_s}{M_s} \vec{W}_{\mu\nu} \vec{W}^{\mu\nu} S^0, \quad (38)$$

$$x_s = - \frac{g_2^2 \tilde{N}_c}{(4\pi)^2} \left(\frac{2f_f I(I+1)(2I+1)M_s}{9M_f} \right). \quad (39)$$

The S^0 exchange between a pair of Higgs doublets and a W pair then generates $\hat{\mathcal{O}}_{UW}$ with a coupling (compare (23))

$$\tilde{d} = - \frac{f_f f_\phi \tilde{N}_c}{18\pi^2} \left(\frac{M_W^2}{M_f M_s} \right) I(I+1)(2I+1). \quad (40)$$

Note that the size of \tilde{d} depends on the unknown factor $f_f f_\phi$. Comparing to the \mathcal{O}_W case, we remark that if this factor is of the order of 1, or even of the electroweak order g_2^2 , then we expect that $|\tilde{d}/\lambda_W| \gtrsim 30$; (compare (35,40)).

As we see from the expressions for the various coupling constants, the above models illustrate how NP can generate a possibly strong selection among the induced operators. This selection can appear as a result of the chosen spectrum in NP and of the precise dynamics determining the size of the various couplings. It may even take the aspect of a symmetry. Whether it really corresponds to a basic symmetry of NP is an open question. This way, NP can prohibit the generation of unwanted large $\Delta\rho$ or $W - B$ mixing effects, ruled out by precise tests.

We also want to make several remarks about the size of the effects. In models A and B' we can naturally accommodate larger couplings for the Higgs involving operators, than for those involving W fields alone; *i.e.* ($\mathcal{O}_{\Phi,2,3} \gtrsim \mathcal{O}_{UW} \gtrsim \mathcal{O}_W$). Of course, this is only true provided scalar elementary NP bosons exist. Otherwise the Higgs involving operators vanish. The couplings of the purely gauge dependent operators are rather

⁶Irrelevant terms like $(S^0)^4$ and $(S^0)^2 \langle \hat{U} \hat{U}^\dagger \rangle$ are omitted in (37), and the vacuum expectation value of S^0 is assumed zero.

weak. This weakness most probably arises from the fact that they solely depend on gauge couplings, for which there is no freedom. On the contrary, the Higgs involving operators naturally share the well known arbitrariness of the scalar sector of SM. And it is this lack of knowledge for the scalar sector, which makes operators like \mathcal{O}_{UW} very interesting. We also repeat that in model A, the $SU(2)_c$ violating coupling $g_{\psi 2}$ is experimentally constrained to be very small.

Turning now to the purely gauge operators above, we note that the $\mathcal{O}_W/\overline{\mathcal{O}}_{DW}$ ratio is very model dependent in sign and magnitude. As mentioned before, the physical contents of these operators is very different. \mathcal{O}_W includes no 2-pt function, and its 3-pt function defines the genuine quadruple gauge boson coupling λ_W [14]. $\overline{\mathcal{O}}_{DW}$, on the other hand, has 2-pt and 3-pt gauge vertices which are just d'Alembertian derivatives of the standard kinetic terms. For this reason the $\overline{\mathcal{O}}_{DW}$ coupling does not seem to bring much NP feature. It is thus conceivable that \mathcal{O}_W and $\overline{\mathcal{O}}_{DW}$ may get very different NP effects.

It is well known that present precision measurements strongly constrain $\overline{\mathcal{O}}_{DW}$ through its 2-pt function part. Do these already exclude our models A and B? We argue below that this is not necessarily the case, since non-perturbative effects due to possible existence of NP resonances could avoid the conclusion that \mathcal{O}_W and $\overline{\mathcal{O}}_{DW}$ have comparable couplings. As an example we consider the possibility of vector bosons V , which are non-relativistic bound states of either the Ψ particles in model A, or the F particles in model B. The gauge boson 2-pt function receives then a contribution from W-V mixing, which involves the V wave function or some of its derivatives at the origin. This effect should have the standard SM structure, and therefore it should be associated to $\overline{\mathcal{O}}_{DW}$. On another hand the quadruple 3-boson coupling described by \mathcal{O}_W receives a contribution involving an integral over the full V wave function. Depending on the relation between the average radius \bar{r} of the V wave function and $1/\Lambda_{NP}$, different situations on the $\mathcal{O}_W/\overline{\mathcal{O}}_{DW}$ ratio may arise. Thus, in case $\bar{r} \gg 1/\Lambda_{NP}$, we would expect that both, the $\overline{\mathcal{O}}_{DW}$ and the \mathcal{O}_W couplings depend on the short distance behaviour of the V wave function. In such a case it may be natural to expect these couplings to be similar. Another more interesting scenario may arise in case $\bar{r} \lesssim 1/\Lambda_{NP}$ since then \mathcal{O}_W depends on the full wave function, which may in turn imply that \mathcal{O}_W becomes very different from $\overline{\mathcal{O}}_{DW}$. Thus, the present strong constraints on $\overline{\mathcal{O}}_{DW}$, should not be taken as excluding the possibility of considerable \mathcal{O}_W contributions in future experiments.

In this paper we have presented dynamical pictures where NP generates at a scale Λ_{NP} a hierarchy of mainly $SU(2)_c$ invariant operators. The basic characteristic of our models is the assumption that the new degrees of freedom determining NP have vanishing hypercharge. The natural leading terms in the induced effective lagrangian involve operators containing Higgs fields; *i.e.* ($\mathcal{O}_{\Phi 2,3}$ and \mathcal{O}_{UW}). Whether \mathcal{O}_W is enhanced with respect to the strongly constrained $\overline{\mathcal{O}}_{DW}$, is a more model dependent question. In model A, an $SU(2)_c$ violating operator is also generated at Λ_{NP} , determined by an independent coupling $g_{\psi 2}$ which is constrained by present data to be very small. On the contrary, in Model B only $SU(2)_c$ conserving operators appear. It is also worth mentioning that if the hypercharge of the Ψ or F NP particles were non vanishing, then the operators \mathcal{O}_{DB} , \mathcal{O}_{BW} and \mathcal{O}_{UB} would also be generated in model A, while in model B only \mathcal{O}_{DB} . The

operators $\mathcal{O}_{W\Phi}$ and $\mathcal{O}_{B\Phi}$ never appear in such models at the level of $dim = 6$ operators. In type A models, $\mathcal{O}_{W\Phi}$ or $\mathcal{O}_{B\Phi}$, multiplied by $\Phi^\dagger\Phi$, first arise at the $dim = 8$ level .

The overall picture is consistent with the view that NP has something to do with the scalar sector. The non observation of some NP effect in the light fermionic sector at LEP1, does not prevent potentially large NP effects to appear in the bosonic sector. These features should be tested at future colliders through gauge boson and Higgs production. They would then provide valuable information on the New Physics properties and its possible origin.

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